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The Minimal Type-I Seesaw Model and Flavor-dependent Leptogenesis

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In this talk, we first give a brief review of the so-called minimal seesaw models and then concentrate on the minimal type-I seesaw model with two almost degenerate right-handed Majorana neutrinos of $\mathcal{O}(1\ \text{TeV})$. A specific texture of the neutrino Yukawa coupling matrix is proposed to achieve the nearly tri-bimaximal neutrino mixing pattern. This ansatz predicts (1) $\theta_{23} = \pi/4$, $|\delta| = \pi/2$ and $\sin^2\theta_{12} = (1 - 2\tan^2\theta_{13})/3$ in the $m_1 = 0$ case; and (2) $\theta_{23} = \pi/4$ and $\theta_{13} = \delta = 0$ in the $m_3 = 0$ case. In both cases, the cosmological baryon number asymmetry can be explained via the resonant leptogenesis mechanism. Finally, we demonstrate the significance of flavor-dependent effects in our leptogenesis scenario.

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Recent neutrino oscillation experiments have convinced us that neutrinos are massive and lepton flavors are mixed.[1] In order to accommodate tiny neutrino masses, one may immediately extend the standard model by introducing three right-handed Majorana neutrinos. The gauge-invariant Lagrangian relevant to lepton masses reads

$$-\mathcal{L}_{\rm mass} = Y_l \overline{l_{\rm L}} H E_{\rm R} + Y_\nu \overline{l_{\rm L}} \tilde{H} N_{\rm R} + \frac{1}{2}\overline{N^c_{\rm R}} M_{\rm R} N_{\rm R} + {\rm h.c.}\ . \tag{1}$$

After the spontaneous gauge symmetry breaking, the Dirac mass matrices of charged leptons and neutrinos are given by $M_l = Y_l v$ and $M_{\rm D} = Y_\nu v$, where $v = \langle H \rangle \approx 174$ GeV is the vacuum expectation value of the neutral Higgs field. The effective mass matrix of three light neutrinos arises from the well-known seesaw formula $M_\nu \approx -M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T$. The lightness of left-handed Majorana neutrinos is therefore attributed to the heaviness of right-handed Majorana neutrinos, while the phenomenon of flavor mixing is due to the mismatch between the diagonalizations of M_l and M_ν. One appealing advantage of the seesaw models is the realization of baryogenesis via leptogenesis: the lepton number asymmetries are first generated from the CP-violating and out-of-equilibrium decays of heavy Majorana neutrinos and then converted into the baryon number asymmetry by means of the $(B-L)$-conserving sphaleron interaction.[2] However, a seesaw model is usually plagued with too many free parameters even in the basis where the mass matrices

of charged leptons and right-handed neutrinos are simultaneously diagonal. Let us explicitly count the number of model parameters. Besides the right-handed Majorana neutrino masses M_i (for $i = 1, 2, 3$), there remain fifteen free parameters in the Dirac neutrino mass matrix $M_{\rm D}$ after three unphysical phases are removed by redefining the charged-lepton fields. At low energies, there are only nine observables: three light neutrino masses m_i (for $i = 1, 2, 3$), three neutrino mixing angles θ_{ij} (for $ij = 12, 23, 13$), and three CP-violating phases (δ, ρ, σ). Hence a generic seesaw model cannot make any specific predictions, unless some additional assumptions are imposed on it.

The minimal type-I seesaw model includes only two heavy right-handed Majorana neutrinos.[3,4] In this case, $M_{\rm D}$ is a 3×2 complex matrix which in general has nine physical parameters. Hence there are totally eleven parameters in this simplified seesaw model. In contrast, there are totally eighteen parameters in the conventional seesaw model with three right-handed Majorana neutrinos. An intrinsic feature of the minimal type-I seesaw model is that the lightest neutrino must be massless and only one of the Majorana CP-violating phases (ρ or σ) survives. It is well known that this is the most economical seesaw scenario that can interpret both the neutrino masses and the cosmological baryon number asymmetry. Alternatively, the introduction of a heavy triplet scalar into the standard model can also give rise to tiny Majorana masses of three known neutrinos.[5] From the viewpoint of grand unified theories, however, the most natural choice is to introduce both the right-handed neutrino singlets and the scalar triplet. It is easy to show that only one triplet Higgs and one right-handed Majorana neutrino are enough to account for both the neutrino oscillation experiments and the observed matter-antimatter asymmetry of the universe. This scenario is usually referred to as the minimal type-II seesaw model.[6]

We now propose an intriguing minimal type-I seesaw model with two nearly degenerate right-handed Majorana neutrinos of $\mathcal{O}(1~{\rm TeV})$, which may be accessible at the forthcoming Large Hadron Collider. Let us begin with a useful parametrization:[7]

$$M_{\rm D}^{(1)} = V_0 \begin{pmatrix} 0 & 0 \\ x & 0 \\ 0 & y \end{pmatrix} U \; , \qquad M_{\rm D}^{(3)} = V_0 \begin{pmatrix} x & 0 \\ 0 & y \\ 0 & 0 \end{pmatrix} U \tag{2}$$

for $m_1 = 0$ and $m_3 = 0$ cases, where V_0 and U are 3×3 and 2×2 unitary matrices, respectively. Then the seesaw relation $M_\nu = -M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T$ implies that the neutrino mixing depends primarily on V_0 and the decays of heavy neutrinos rely mainly on U. Hence we take V_0 to be the tri-bimaximal mixing pattern[8]

$$V_0 = \begin{pmatrix} 2/\sqrt{6} & 1/\sqrt{3} & 0 \\ -1/\sqrt{6} & 1/\sqrt{3} & 1/\sqrt{2} \\ 1/\sqrt{6} & -1/\sqrt{3} & 1/\sqrt{2} \end{pmatrix} \; , \tag{3}$$

which is compatible very well with the best fit of current experimental data.[1] On

the other hand, the unitary matrix U can be parameterized as

$$U = \begin{pmatrix} \cos\vartheta & \sin\vartheta \\ -\sin\vartheta & \cos\vartheta \end{pmatrix} \begin{pmatrix} e^{-i\alpha} & 0 \\ 0 & e^{+i\alpha} \end{pmatrix} . \qquad (4)$$

Since α is the only phase parameter in our model, it should be responsible both for the CP violation in neutrino oscillations and for the CP violation in N_i decays. For simplicity, here we fix $\vartheta = \pi/4$ and highlight the role of α in neutrino mixing and leptogenesis. In order to implement the idea of resonant leptogenesis, we suppose that two heavy Majorana neutrino masses are highly degenerate; i.e., the magnitude of $r \equiv (M_2 - M_1)/M_2$ is strongly suppressed.[9]

Given $|r| < \mathcal{O}(10^{-4})$, the explicit form of M_ν can reliably be formulated from the seesaw relation $M_\nu = -M_{\rm D} M_{\rm R}^{-1} M_{\rm D}^T$ by neglecting the tiny mass splitting between N_1 and N_2. In this good approximation, we diagonalize M_ν through $V^\dagger M_\nu V^* = {\rm Diag}\,\{m_1, m_2, m_3\}$, where V is just the neutrino mixing matrix. In the $m_1 = 0$ case, three mixing angles and the Dirac CP-violating phase are determined by

$$\sin^2\theta_{12} = \frac{1-\sin^2\theta}{3-\sin^2\theta} \ , \quad \theta_{23} = \frac{\pi}{4} \ , \quad \sin^2\theta_{13} = \frac{\sin^2\theta}{3} \ , \quad \delta = -\frac{\pi}{2} \ , \qquad (5)$$

where θ is given by $\tan 2\theta = 2\omega \tan 2\alpha/(1+\omega^2)$ with $\omega \equiv x/y \in (0,1)$. It is straightforward to derive $\sin^2\theta_{12} = (1 - 2\tan^2\theta_{13})/3$. Taking account of $m_2 = \sqrt{\Delta m^2_{21}}$ and $m_3 = \sqrt{\Delta m^2_{21} + |\Delta m^2_{32}|}$, we arrive at $m_2 \approx 8.9 \times 10^{-3}$ eV and $m_3 \approx 5.1 \times 10^{-2}$ eV by using $\Delta m^2_{21} \approx 8.0 \times 10^{-5}$ eV^2 and $|\Delta m^2_{32}| \approx 2.5 \times 10^{-3}$ eV^2 as the typical input.[1] The values of m_2 and m_3, together with $\theta_{13} < 10^\circ$, lead to $0.39 \lesssim \omega \lesssim 0.42$, $0^\circ \lesssim \alpha \lesssim 23^\circ$ and $0^\circ \lesssim \theta \lesssim 18^\circ$. We find that $m_2 \approx x^2/M_2$ and $m_3 \approx y^2/M_2$ are good approximations for $\alpha \lesssim 10^\circ$. For the inverted neutrino mass hierarchy ($m_3 = 0$), we have

$$\sin^2\theta_{12} = \frac{1+\sin^2\theta}{3} \ , \qquad \theta_{23} = \frac{\pi}{4} \ , \qquad \theta_{13} = 0 \ , \qquad \delta = 0 \ . \qquad (6)$$

Taking account of $m_1 = \sqrt{|\Delta m^2_{32}| - \Delta m^2_{21}}$ and $m_2 = \sqrt{|\Delta m^2_{32}|}$, we get $m_1 \approx 4.9 \times 10^{-2}$ eV and $m_2 \approx 5.0 \times 10^{-2}$ eV by inputting $\Delta m^2_{21} \approx 8.0 \times 10^{-5}$ eV^2 and $|\Delta m^2_{32}| \approx 2.5 \times 10^{-3}$ eV^2. Given $30^\circ < \theta_{12} < 38^\circ$,[1] θ is found to lie in the range $0 \lesssim \theta \lesssim 22^\circ$. Furthermore, the values of m_1 and m_2 allow us to get $0^\circ \lesssim \alpha \lesssim 22^\circ$ and $0.991 \lesssim \omega \lesssim 0.992$. The neutrino masses reliably approximate to $m_1 \approx x^2/M_2$ and $m_2 \approx y^2/M_2$ for $\alpha \lesssim 10^\circ$ in this case. From a phenomenological point of view, the scenario with $m_1 = 0$ is more favored and more interesting than the scenario with $m_3 = 0$. Both of them can be tested in the near future.

We proceed to calculate the cosmological baryon number asymmetry via the flavor-dependent resonant leptogenesis. Note that all the Yukawa interactions of charged leptons are in thermal equilibrium at the TeV scale, so the lepton number asymmetry for each flavor should be treated separately due to its distinct thermal history.[10] In the framework of resonant leptogenesis, it is straightforward to compute the CP-violating asymmetry between $N_i \to l_\alpha + H^{\rm c}$ and $N_i \to l^{\rm c}_\alpha + H$ decays

for each lepton flavor α ($= e$, μ or τ):

$$\varepsilon_{i\alpha} = \frac{8\pi\left(M_i^2 - M_j^2\right) \mathrm{Im}\left\{\left(Y_\nu\right)_{\alpha j}\left(Y_\nu\right)^*_{\alpha i} M_i\left[M_i\left(Y_\nu^\dagger Y_\nu\right)_{ji} + M_j\left(Y_\nu^\dagger Y_\nu\right)_{ij}\right]\right\}}{\left[64\pi^2\left(M_i^2 - M_j^2\right)^2 + M_i^4\left(Y_\nu^\dagger Y_\nu\right)^2_{jj}\right]\left(Y_\nu^\dagger Y_\nu\right)_{ii}} , \quad (7)$$

where i and j run over 1 and 2 but $i \neq j$. To take account of the inverse decays and lepton-number-violating scattering processes, we define the corresponding decay parameters $K_{i\alpha} \equiv \left|(Y_\nu)_{\alpha i}\right|^2 K_i / \left(Y_\nu^\dagger Y_\nu\right)_{ii}$, where $K_i \equiv \Gamma_i/H$ at $T = M_i$ with $\Gamma_i = \left(Y_\nu^\dagger Y_\nu\right)_{ii} M_i/(8\pi)$ being the total decay width of N_i and $H(T) = 1.66\sqrt{g_*}T^2/M_{\rm planck}$ being the Hubble parameter. Note that these quantities can be explicitly figured out with the help of Eqs. (2), (3) and (4) as well as the constraints from neutrino masses and mixing parameters. In the strong washout regime $K_{i\alpha} \gtrsim 1$, the survival of lepton number asymmetries is approximately characterized by the efficiency factors[11]

$$\kappa_{i\alpha} \approx \frac{2}{K_\alpha z_{\rm B}(K_\alpha)}\left[1 - \exp\left(-\frac{K_\alpha z_{\rm B}(K_\alpha)}{2}\right)\right] , \quad (8)$$

with $z_{\rm B}(K_\alpha) \simeq 2 + 4K_\alpha^{0.13}\exp\left(-2.5/K_\alpha\right)$ and $K_\alpha = \sum_i K_{i\alpha}$. The final baryon number asymmetry can be estimated as $\eta^{\rm f}_{\rm B} \approx -0.96 \times 10^{-2}\sum_i\sum_\alpha\left(\varepsilon_{i\alpha}\kappa_{i\alpha}\right)$. It is found that the observed baryon number asymmetry $\eta \approx 6.1 \times 10^{-10}$ can indeed be achieved.[7] For comparison, we also examine the final baryon asymmetry in the one-flavor approximation. Summing the CP-violating asymmetries over all flavors $\varepsilon_i = \sum_\alpha \varepsilon_{i\alpha}$ and using the flavor-independent efficiency factors $\kappa_i \approx 0.5/\left(\sum_i K_i\right)^{1.2}$, one gets the final baryon asymmetry $\eta_{\rm B} \approx -0.96 \times 10^{-2}\sum_i\left(\varepsilon_i\kappa_i\right)$. In our scenario,

$$\frac{\eta^{\rm f}_{\rm B}}{\eta_{\rm B}} = \frac{\sum_{i,\alpha}\varepsilon_{i\alpha}\kappa_{i\alpha}}{\sum_i\varepsilon_i\kappa_i} \approx \begin{cases} -0.71 & (m_1 = 0) , \\ +52.1 & (m_3 = 0) . \end{cases} \quad (9)$$

We see that the flavor-independent prediction for $\eta_{\rm B}$ becomes negative in the $m_1 = 0$ case, while it is enhanced by a factor ~ 50 in the $m_3 = 0$ case.

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